$B \to \chi_{c1}(1P, 2P)K$ decays in QCD factorization and $X(3872)$

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$B \to \chi_{c1}(1P, 2P)K$ decays are studied in QCD factorization by treating charmonia as nonrelativistic bound states. No infrared divergences exist in the vertex corrections, while the logarithmic end-point singularities in the hard spectator corrections can be regularized by a momentum cutoff. Within certain uncertainties, we find that the $B \to \chi_{c1}(2P)K$ decay rate can be comparable to $B \to \chi_{c1}(1P)K$ and get $\mathrm{Br}(B^0 \to \chi'_{c1}K^0) = \mathrm{Br}(B^+ \to \chi'_{c1}K^+) \approx (2\text{–}4) \times 10^{-4}$. This might imply a possible interpretation for the newly discovered X(3872): that this state has a dominant $J^{PC} = 1^{++}(2P)$ $c\bar{c}$ component, but it is mixed with a substantial $D^0\bar{D}^{*0} + D^{*0}\bar{D}^0$ component.

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The naively factorizable decay [1] $B \to \chi_{c1}K$ was studied [2] in the QCD factorization approach [3,4], in which the nonfactorizable vertex and spectator corrections were also estimated, but the numerical results were four times smaller than experimental data. Recently, these decays were also studied in the PQCD approach [5]. In both the above approaches, light-cone distribution amplitudes (LCDAs) were used to describe χ_{c1}. As argued in Ref. [6], a more appropriate description of charmonium is to use the nonrelativistic (NR) wave functions which can be expanded in terms of the relative momentum q between charm and anticharm quarks. This argument is based on the nonrelativistic nature of heavy quarkonium [7]. With careful studies, we find that the two descriptions (i.e., LCDAs and NR) are equivalent for the S-wave charmonium states (see, e.g., Ref. [8]), but in the case of P-wave states the light-cone descriptions lose some important contributions in the leading twist approximation. This is not surprising, since q can be neglected in the S-wave states, but it cannot be neglected for the P-wave states, even in the leading order approximation.

On the phenomenological hand, the study of $B \to \chi_{c1}(2P)K$ may help clarify the nature of the recently discovered resonance $X(3872)$ [9], since the measurements for $X(3872)$ favor $J^{PC} = 1^{++}$ [10], and hence $\chi_{c1}(2P)$ becomes one of the possible assignments for it. On the other hand, aside from the conventional charmonium [11,12], a loosely bound S-wave molecule of $D^0\bar{D}^{*0} + D^{*0}\bar{D}^0$ has been suggested for $X(3872)$ [13,14].

Motivated by the above considerations, in this paper we study the decays $B \to \chi_{c1}(1P, 2P)K$ within the framework of QCD factorization by treating the charmonia $\chi_{c1}(1P, 2P)$ as nonrelativistic bound states, with m_c/m_b taken to be a fixed value in the heavy b-quark limit. We will estimate the production rate of $\chi_{c1}(2P)$ and argue that the $X(3872)$ may be dominated by the $\chi_{c1}(2P)$ charmonium but mixed with some $D^0\bar{D}^{*0} + D^{*0}\bar{D}^0$ component.

In the nonrelativistic bound state picture, charmonium can be described by the color singlet NR wave function. Let p be the total momentum of the charmonium and $2q$ be the relative momentum between c and $\bar{c}$ quarks; then $v^2 \sim 4q^2/p^2 \sim 0.25$ can be treated as a small expansion parameter [7]. For P-wave charmonium χ_{c1}, because the wave function at the origin $\mathcal{R}_P(0) = 0$, which corresponds to the zeroth order in q, we must expand the amplitude to the first order in q. Thus, we have

$$\begin{aligned}\mathcal{M}(B \to \chi_{c1}K) &= \sum_{L_z,S_z} \langle 1L_z; 1S_z | 1J_z \rangle \int \frac{d^4q}{(2\pi)^3} q_\alpha \delta(q^0) \psi^*_{1M}(q) \\ &\quad \times \mathrm{Tr}[\mathcal{O}^\alpha(0) P_{1S_z}(p,0) + \mathcal{O}(0) P^\alpha_{1S_z}(p,0)], \qquad (1)\end{aligned}$$

where $\mathcal{O}(q)$ represent the rest of the decay amplitudes, $P_{1S_z}(p,q)$ is the spin-triplet projection operator, and $\mathcal{O}^\alpha$, P^α stand for the derivatives of $\mathcal{O}$, P with respect to the relative momentum q_α [6]. The amplitudes $\mathcal{O}(q)$ can be further factorized as the product of $B \to K$ form factors and a hard kernel, or as the convolution of a hard kernel with light-cone wave functions of B mesons and K mesons, within the QCD factorization approach.

After q^0 is integrated out, the integral in Eq. (1) is proportional to the derivative of the P-wave wave function at the origin by

$$\int \frac{d^3q}{(2\pi)^3} q^\alpha \psi^*_{1M}(q) = -i\varepsilon^{*\alpha}(L_z)\sqrt{\frac{3}{4\pi}}\mathcal{R}'_P(0), \qquad (2)$$

where $\varepsilon^\alpha(L_z)$ is the polarization vector of an angular-momentum-1 system, and the value of $\mathcal{R}'_P(0)$ for charmonia can be found in, e.g., Ref. [15].

In contrast to the NR description of χ_{c1}, the K meson is described by LCDAs [3]:

$$\langle K(p')|\bar{s}_\beta(z_2)d_\alpha(z_1)|0\rangle = \frac{if_K}{4}\int_0^1 dx e^{i(yp'\cdot z_2+\bar{y}p'\cdot z_1)}\{\not{p}'\gamma_5\phi_K(y)\}_{\alpha\beta}, \quad (3)$$

where y and $\bar{y} = 1 - y$ are the momentum fractions of the s and $\bar{d}$ quarks inside the K meson, respectively, and $\phi_K(x)$ is the leading twist LCDA of the K meson. The masses of light quarks and K mesons are neglected in the heavy quark limit.

The effective Hamiltonian for $B \to \chi_{c1}K$ reads [16]

$$\mathcal{H}_{\rm eff} = \frac{G_F}{\sqrt{2}}\left(V_{cb}V_{cs}^*(C_1\mathcal{O}_1 + C_2\mathcal{O}_2) - V_{tb}V_{ts}^*\sum_{i=3}^{6}C_i\mathcal{O}_i\right), \quad (4)$$

where G_F is the Fermi constant, C_i are the Wilson coefficients, and $V_{q_1q_2}$ are the CKM matrix elements. The relevant four-fermion operators $\mathcal{O}_i$ can be found in Ref. [6].

According to Refs. [3,4], all nonfactorizable corrections are due to the interactions shown in Fig. 1. These corrections, with operators $\mathcal{O}_i$ inserted, contribute to the amplitude $\mathcal{O}(q)$ in Eq. (1), where the external lines of charm and anticharm quarks have been truncated. Taking nonfactorizable corrections in Fig. 1 into account, the decay amplitude for $B \to \chi_{c1}K$ in QCD factorization is written as

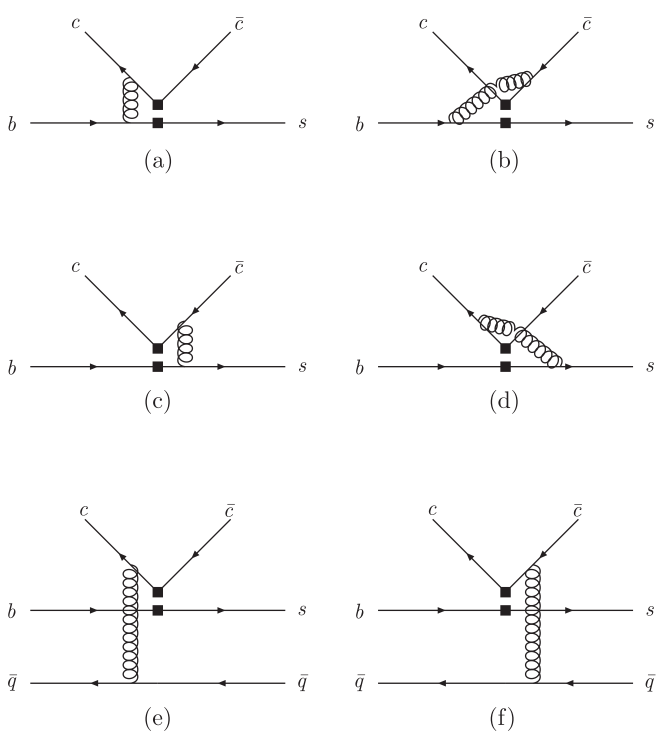

FIG. 1. Feynman diagrams for vertex and spectator corrections to $B \to \chi_{c1}K$.

$$i\mathcal{M} = \frac{G_F}{\sqrt{2}}[V_{cb}V_{cs}^*a_2 - V_{tb}V_{ts}^*(a_3 - a_5)] \times 12i\sqrt{\frac{2}{\pi M}}\mathcal{R}_P'(0)\epsilon^*\cdot p_B F_1(M^2), \quad (5)$$

where ϵ is the polarization vector of χ_{c1} [17]. Here F_1 is the $B \to K$ form factor, and we have used the approximate relation [18]

$$F_0(M^2)/F_1(M^2) \cong 1 - z, \quad (6)$$

with $z = M^2/m_B^2 \approx 4m_c^2/m_b^2$ and M as the mass of χ_{c1}, to simplify the structure of Eq. (5).

The coefficients a_i ($i = 2, 3, 5$) in the naive dimension regularization scheme are given by

$$\begin{aligned} a_2 &= C_2 + \frac{C_1}{N_c} + \frac{\alpha_s}{4\pi}\frac{C_F}{N_c}C_1\left(-18 + 12\ln\frac{m_b}{\mu} + f_I + f_{II}\right),\\ a_3 &= C_3 + \frac{C_4}{N_c} + \frac{\alpha_s}{4\pi}\frac{C_F}{N_c}C_4\left(-18 + 12\ln\frac{m_b}{\mu} + f_I + f_{II}\right),\\ a_5 &= C_5 + \frac{C_6}{N_c} - \frac{\alpha_s}{4\pi}\frac{C_F}{N_c}C_6\left(-6 + 12\ln\frac{m_b}{\mu} + f_I + f_{II}\right), \end{aligned} \quad (7)$$

where $C_F = (N_c^2 - 1)/(2N_c)$, and μ is the QCD renormalization scale.

The function f_I is calculated from the four vertex correction diagrams (a, b, c, d) in Fig. 1 and reads

$$f_I = \frac{2z}{2-z} - \frac{4z\log(4)}{2-z} - \frac{4z^2\log(z)}{(1-z)(2-z)} + \frac{4(3-2z)(1-z)(\log(1-z) - i\pi)}{(2-z)^2}. \quad (8)$$

We find that the infrared divergences are canceled between diagrams (a) and (b), and between (c) and (d) in Fig. 1, and this cancellation is independent of whether the relation in Eq. (6) is used. On the other hand, this function is different from that in Eq. (11) of Ref. [2] even when a nonrelativistic limit wave function $\phi_{\chi_{c1}}^{\rm NR}(u) = \delta(u - 1/2)$ is adopted, as we have mentioned.

For the two spectator correction diagrams (e, f) in Fig. 1, the off-shell-ness of the gluon is naturally associated with a scale $\mu_h \sim \sqrt{m_b\Lambda_{\rm QCD}}$, rather than $\mu_h \sim m_b$. Following Refs. [3,4], we choose $\mu = \sqrt{m_b\Lambda_h} \approx 1.4$ GeV with $\Lambda_h = 0.5$ GeV in calculating the hard spectator function f_{II}, and then, in the leading twist approximation, we get

$$f_{II} = \frac{\alpha_s(\mu_h)C_i(\mu_h)}{\alpha_s(\mu)C_i(\mu)}\frac{8\pi^2}{N_c}\frac{f_Kf_B}{F_1(M^2)m_B^2}\frac{1}{1-z}\int_0^1 d\xi\frac{\phi_B(\xi)}{\xi} \times \int_0^1 dy\frac{\phi_K(y)}{y}\left[1 + \frac{z}{y(1-z)}\right], \quad (9)$$

where ξ is the momentum fraction of the spectator quark in the B meson, and $C_i(\mu_h)$ ($i = 1, 4, 6$) are the NLO Wilson coefficients which can be evaluated by the renormalization group approach [16].

The spectator contribution depends on the wave function ϕ_B through the integral

$$\int_0^1 d\xi \frac{\phi_B(\xi)}{\xi} \equiv \frac{m_B}{\lambda_B}. \tag{10}$$

Since $\phi_B(\xi)$ is appreciable only for ξ of order $\Lambda_{\rm QCD}/m_B$, λ_B is of order $\Lambda_{\rm QCD}$. We will choose $\lambda_B \approx 300$ MeV in the numerical calculations [4].

If we choose the asymptotic form of the K-meson twist-2 LCDA, $\phi_K(y) = 6y(1-y)$, we can find logarithmic end-point singularities in Eq. (9) just like that in Ref. [2], and we parameterize it in a simple way [4]:

$$\int \frac{dy}{y} = \ln \frac{m_B}{\Lambda_h} \approx 2.4. \tag{11}$$

The mass of $\chi_{c1}(1P)$, $M_{\chi_{c1}} = 3.511$ GeV, is known, but the mass of the missing charmonium $\chi_{c1}(2P)$ has to be estimated by, say, potential models. We choose $M_{\chi'_{c1}} = 3.953$ GeV, following Ref. [19]. Then the form factor $F_1(M^2)$ can be determined by light-cone sum rules [20]:

$$F_1\left(M^2_{\chi_{c1}}\right) = 0.80, \qquad F_1\left(M^2_{\chi'_{c1}}\right) = 1.14. \tag{12}$$

We also choose $M_{\chi'_{c1}} = 3.872$ GeV and $F_1(M^2_{\chi'_{c1}}) = 1.06$ to study if the $X(3872)$ behaves like a $\chi_{c1}(2P)$ in its b-production process.

For numerical analysis, we use the following input parameters [21]:

$$\begin{aligned}
&m_b = 4.8 \text{ GeV}, \qquad m_B = 5.28 \text{ GeV} \\
&f_K = 160 \text{ MeV}, \qquad f_B = 216 \text{ MeV}, \\
&\mathcal{R}'_{1P}(0) = \mathcal{R}'_{2P}(0) = \sqrt{0.1} \text{ GeV}^{5/2}, \\
&C_1(\mu) = 1.21\,(1.082), \qquad C_2(\mu) = -0.40\,(-0.185), \\
&C_3(\mu) = 0.03\,(0.014), \qquad C_4(\mu) = -0.05\,(-0.035), \\
&C_5(\mu) = 0.01\,(0.009), \qquad C_6(\mu) = -0.07\,(-0.041), \\
&\alpha_s(\mu) = 0.35\,(0.22).
\end{aligned} \tag{13}$$

In Eq. (13), the μ-dependent quantities at $\mu_h = 1.4$ GeV ($\mu = 4.4$ GeV) are shown without (with) parentheses.

Using the above inputs, we get the results of the coefficients a_i which are listed in Table I. With the help of these coefficients a_i, we calculate the decay branching ratios of decays $B \to \chi_{c1}(1P, 2P)K$ with two different choices of $M_{\chi'_{c1}}$ and get

$$\begin{aligned}
&\mathrm{Br}(B^0 \to \chi_{c1}(3511)K^0) = 1.79 \times 10^{-4}, \\
&\mathrm{Br}(B^0 \to \chi'_{c1}(3953)K^0) = 1.81 \times 10^{-4}, \\
&\mathrm{Br}(B^0 \to \chi'_{c1}(3872)K^0) = 1.78 \times 10^{-4}.
\end{aligned} \tag{14}$$

TABLE I. The coefficients a_i of $B \to \chi_{c1}(1P, 2P)K$ with different choices of $M_{\chi'_{c1}}$.

	a_2	a_3	a_5
$\chi_{c1}(3511)$	$0.199 - 0.051i$	$0.000 + 0.002i$	$0.004 - 0.002i$
$\chi'_{c1}(3953)$	$0.247 - 0.042i$	$-0.002 + 0.001i$	$0.007 - 0.002i$
$\chi'_{c1}(3872)$	$0.236 - 0.044i$	$-0.002 + 0.001i$	$0.006 - 0.002i$

Our prediction of $\mathrm{Br}(B^0 \to \chi_{c1}(3511)K^0)$ is about two times larger than that in Ref. [2], although it is still about two times smaller than the recent data [22]. The difference between the theoretical predictions and experimental data may not be as serious as it looks if we take into account the following uncertainties: (i) We have used a moderate value of $\mathcal{R}'_{1P}(0)$ predicted by different potential models [15] in our calculation, and a larger value of $\mathcal{R}'_{1P}(0)$ may enhance our prediction in Eq. (13) significantly. (ii) In our evaluation of f_{II}, we only use the leading twist LCDAs of the K meson, and large uncertainties will arise from the chirally enhanced higher twist effects [18]. (iii) Since the squared velocity v^2 of the charm quark in charmonium is about 0.25–0.30, the relativistic corrections may be important for these decays.

Note that although the form factor in Eq. (12) and the coefficient a_2 in Table I increase evidently as the charmonium mass increases, the decreased phase space and kinematic factors in Eq. (5) will create a balance and result in similar decay branching ratios in the charmonium mass region 3.51–3.95 GeV, as shown in Eq. (14). If we neglect the order-α_s corrections (i.e., in the naive factorization [1]), the ratios among these three branching fractions in Eq. (14) would become 1:0.74:0.69. As a rough estimate, we expect the branching ratios for $\chi_{c1}(2P)$ to be

$$\mathcal{B}^0 \equiv \mathrm{Br}(B^0 \to \chi'_{c1}K^0) \approx (2\text{–}4) \times 10^{-4}, \tag{15}$$

where the values are taken between the calculated values for $\chi_{c1}(2P)$ and the experimental values for $\chi_{c1}(1P)$.

In fact, the end-point singularities in Eq. (9) make the factorization break down even at the leading twist level, which indicates that the soft spectator interactions may be important. Similar logarithmic end-point singularities also emerge in the twist-3 spectator interactions in B decay to two light mesons, where the soft contributions, which are parameterized by the momentum cutoff $\Lambda_h \sim 500$ MeV, are numerically important [4]. Here we estimate the soft spectator contributions in Eq. (11) in a similar way to that in Ref. [4], and the numerical results in Eq. (14) suffer from large uncertainties in the momentum cutoff. However, the sensitivity to the cutoff could be canceled to a large extent in the ratio of the production rates of $\chi_{c1}(2P)$ and $\chi_{c1}(1P)$ in B-meson decays. In this sense, the estimation in Eq. (15) should be reasonable.

Moreover, because isospin is conserved in the heavy quark limit, the branching ratio of the charged channel

$\mathcal{B}^{\pm} \equiv \mathrm{Br}(B^{\pm} \to \chi'_{c1} K^0)$ should approximate to the neutral one, and we then predict the ratio

$$R_{\chi'_{c1}} \equiv \frac{\mathcal{B}^0}{\mathcal{B}^{\pm}} \approx \frac{\tau[B^0]}{\tau[B^+]} \approx 0.9, \tag{16}$$

where $\tau[B^0(B^+)]$ is the lifetime of the $B^0(B^+)$ meson.

Comparing Eq. (15) with the measured channel of the $X(3872)$ [9],

$$\begin{aligned} \mathrm{Br}(B^+ \to XK^+) \times \mathcal{B}_X = (1.3 \pm 0.3) \times 10^{-5}, \\ \mathcal{B}_X \equiv \mathrm{Br}(X \to J/\psi \pi^+ \pi^-), \end{aligned} \tag{17}$$

we see that the produced $X(3872)$ looks like the $\chi_{c1}(2P)$ if $\mathcal{B}_X$ is sufficiently small—say, 3%–7%. A similar conclusion has been obtained in a comprehensive analysis of $X(3872)$ production at the Tevatron and B factories [23]. On the other hand, if $X(3872)$ is a loosely bound S-wave molecule of $D^0\bar{D}^{*0}/D^{*0}\bar{D}^0$ [13,24], a model calculation gives a smaller rate [14] compared with Eq. (15):

$$\mathrm{Br}\,(B^+ \to XK^+) = (0.07\text{–}1) \times 10^{-4}, \tag{18}$$

which requires a $\mathcal{B}_X > 10\%$ in order to be consistent with experimental data [Eq. (16)]. They also predict

$$\mathrm{Br}\,(B^0 \to X(3872)K^0) < 0.1\mathrm{Br}(B^+ \to X(3872)K^+). \tag{19}$$

So the measurement of $\mathcal{B}_X$ and $\mathrm{Br}(B^0 \to X(3872)K^0)$ is very helpful to identify the nature of $X(3872)$.

Recently, a preliminary result for a new decay mode, $X \to D^0\bar{D}^0\pi^0$, was found by Belle [25]:

$$\mathrm{Br}(B \to XK) \times \mathrm{Br}(X \to D^0\bar{D}^0\pi^0) = (2.2 \pm 0.7 \pm 0.4) \times 10^{-4}. \tag{20}$$

Equation (20) implies that $\mathcal{B}_X < 10\%$, if it can be confirmed by further measurements. This would disfavor the suggestion that the $X(3872)$ is a loosely bound S-wave molecule of $D^0\bar{D}^{*0}/D^{*0}\bar{D}^0$ with predictions of both decay [24] and production [14].

The above discussions about the $X(3872)$ are based on the assumption that the $X(3872)$ is a pure charmonium $\chi_{c1}(2P)$ state. But this cannot be the case due to the coupled channel effects and $X(3872)$ being in extremely close proximity to the $D^0\bar{D}^{*0}/D^{*0}\bar{D}^0$ threshold. Perhaps a more realistic model for the $X(3872)$ (for further discussions, see Ref. [26]) is that the $X(3872)$ has a dominant $J^{PC} = 1^{++}(2P)$ $c\bar{c}$ component which is mixed with a substantial $D^0\bar{D}^{*0}/D^{*0}\bar{D}^0$ continuum component. (The $D^+\bar{D}^{*-}/D^{*-}\bar{D}^+$ continuum component is kinematically forbidden to be mixed in $X(3872)$, and it is the $u - d$ quark mass difference that causes this isospin violation.) Thus, $X(3872)$ will have the following features: (i) The production of $X(3872)$ in B-meson decays is mainly due to the $J^{PC} = 1^{++}(2P)$ $c\bar{c}$ component, as discussed above. The production of $X(3872)$ at the Tevatron is also due to this $c\bar{c}$ component and associated higher Fock states containing the color-octet $c\bar{c}$ pair and soft gluons. As it was argued [12] for the prompt charmonium production that cross sections of D-wave charmonia [which were suggested as a tentative candidates for $X(3872)$ in Ref. [12]] could be as large as J/ψ or $\psi(2S)$ due to the color octet mechanism, the P-wave $(2P)$ charmonium could also have a comparable production rate to J/ψ or $\psi(2S)$. But this does not seem to be obvious for a loosely bound S-wave molecule of $D^0\bar{D}^{*0}/D^{*0}\bar{D}^0$. (ii) On the other hand, the $D^0\bar{D}^{*0}/D^{*0}\bar{D}^0$ component in $X(3872)$ will be mainly in charge of the hadronic decays of $X(3872)$ into $D^0\bar{D}^{*0}/D^{*0}\bar{D}^0$ or $D^0\bar{D}^0\pi^0$ as well as $J/\psi\rho^0$ and $J/\psi\omega$. The latter two decay modes ($J/\psi\rho^0$ and $J/\psi\omega$) may come from the first decay mode $D^0\bar{D}^{*0}/D^{*0}\bar{D}^0$ and a subsequent rescattering final-state interaction, and therefore have the same decay amplitudes $[\mathrm{A}(J/\psi\rho^0) = \mathrm{A}(J/\psi\omega)]$ that are smaller than the first decay mode amplitude. (iii) A substantial $D^0\bar{D}^{*0}/D^{*0}\bar{D}^0$ component in $X(3872)$ may reduce the production rates in Eq. (15), and will also reduce the $X(3872) \to J/\psi\gamma$ decay width, which can be as small as 11 KeV [11]. (Note that this 2P–1S E1 transition is sensitive to the model details; see, e.g., Ref. [24].) This is much smaller than the hadronic decay widths. But a large rate for $\chi_{c1}(2P) \to \gamma\psi(2S) \approx 60$ keV will be expected. These qualitative features are useful in understanding the nature of $X(3872)$ and should be further tested and studied experimentally and theoretically.

In summary, we study the decays $B \to \chi_{c1}(1P, 2P)K$ in QCD factorization by treating charmonia as nonrelativistic bound states. We find that there are no infrared divergences in the vertex corrections, and the logarithmic end-point singularities from hard spectator interactions can be regularized by a momentum cutoff. Within certain uncertainties, we find that the $B \to \chi_{c1}(2P)K$ decay rate can be comparable to $B \to \chi_{c1}(1P)K$ [in the ratio $\frac{B \to \chi_{c1}(2P)K}{B \to \chi_{c1}(1P)K}$, the uncertainties due to the momentum cutoff in spectator interactions are largely canceled out], and we get $\mathrm{Br}(B^0 \to \chi'_{c1}K^0) = \mathrm{Br}(B^+ \to \chi'_{c1}K^+) \approx (2\text{–}4) \times 10^{-4}$. This might imply that the $X(3872)$ has a dominant $J^{PC} = 1^{++}(2P)$ $c\bar{c}$ component, but it is mixed with a substantial $D^0\bar{D}^{*0} + D^{*0}\bar{D}^0$ component. The qualitative features of $X(3872)$ are discussed and should be further tested and studied.

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Note added.—After this work, measurements by the *BABAR* Collaboration were reported [27]: $\mathrm{Br}(B^+ \to$

$X(3872)K^+) < 3.2 \times 10^{-4}$, $R = \frac{\mathrm{Br}(B^0 \to X(3872)K^0)}{\mathrm{Br}(B^+ \to X(3872)K^+)} = 0.50 \pm 0.30 \pm 0.05$, and a similar conclusion to ours for the $X(3872)$ was also obtained in Ref. [28].

In particular, three years later, the decay $B \to \chi_{c1}K$ was revisited by M. Beneke and L. Vernazza in the Coulombic limit where $m_c v^2 \gg \Lambda_{\mathrm{QCD}}$ [29]. The hard vertex function f_I in Ref. [29] is almost the same as ours in Eq. (8), except for a constant term which can be absorbed in the redefinition of the wave function of χ_{c1}. On the other hand, the end-point singularity of the spectator interaction disappears in the Coulombic limit, which is superseded by the logarithmic dependence of the binding energy of the charmonium [29]. The above way to parameterize the "soft" spectator interaction is similar to that in Eq. (11), where the momentum cutoff Λ_h is used to regularize the singularity, so similar numerical results to ours were obtained in Ref. [29].

As for the phenomenological aspects, more evidence of a substantial $c\bar{c}$ component in $X(3872)$ has been seen by fitting the experimental data for $B \to X(3872)K \to J/\psi \pi^+\pi^- K$ and $B \to X(3872)K \to D^0\bar{D}^{*0}K$. Especially, the obtained branching ratio [30]

$$\mathrm{Br}^{\mathrm{fit}}(B \to X(3872)K) = (3\text{–}5) \times 10^{-4} \qquad (21)$$

is consistent with our prediction in Eq. (15), and this may indicate that the $X(3872)$ is produced in B-meson decays mainly through the $J^{PC} = 1^{++}(2P)$ $c\bar{c}$ component in $X(3872)$, as we have suggested in this paper.

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